Minimal Resonant Leptogenesis and Lepton Flavour Violation

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Abstract. We discuss minimal non-supersymmetric models of resonant leptogenesis, based on an approximate flavour symmetries. As an illustrative example, we consider a resonant τ-leptogenesis model, compatible with universal right-handed neutrino masses at the GUT scale, where the required heavy-neutrino mass splittings are generated radiatively. In particular, we explicitly demonstrate, how a minimum number of three heavy Majorana neutrinos is needed, in order to obtain successful leptogenesis and experimentally testable rates for processes of lepton flavour violation, such as $\mu \to e\gamma$ and $\mu \to e$ conversion in nuclei.

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INTRODUCTION

The observed baryon asymmetry in the Universe (BAU), which amounts to a baryon-to-photon ratio of number densities $\eta_B \approx 6.2 \times 10^{-10}$ [1, 2], provides one of the strongest pieces of evidence for physics beyond the Standard Model (SM) [3]. One interesting scenario for explaining the BAU is leptogenesis [4]. A potentially interesting alternative to GUT-scale leptogenesis is the framework of low-scale resonant leptogenesis (RL) [5, 6]. Within this framework, the lowering of the scale may rely on a dynamical mechanism, in which heavy-neutrino self-energy effects [7] on the leptonic asymmetry become dominant [8] and get resonantly enhanced [5], when a pair of heavy Majorana neutrinos has a mass difference comparable to the heavy neutrino decay widths.

Flavour effects due to heavy-neutrino Yukawa couplings [9] play an important role in models of RL [10, 11]. In particular, in [10] a scenario was put forward, called resonant τ-leptogenesis (RτL), in which the BAU originates from a τ-lepton asymmetry, resonantly produced by quasi-in-equilibrium decays of heavy Majorana neutrinos. In a RτL model, the generated excess in the L_τ number will be converted into the observed BAU, provided the L_τ-violating reactions are not strong enough to wash out such an excess. In such a scenario, the heavy Majorana neutrinos can be as light as 100 GeV and have sizeable couplings to two of the charged leptons, specifically to electrons and muons. Consequently, depending on the flavour dynamics of heavy neutrino Yukawa coupling effects, phenomenologically testable models of RL can be built that could be probed at the LHC or in low energy experiments of lepton number violation (LNV) and lepton flavour violation (LFV), such as $\mu \to e\gamma$ and neutrinoless double beta decay. As we demonstrated in [12], RℓL models can be constructed with a degenerate heavy neutrino mass spectrum at the GUT scale. In a minimal non-supersymmetric framework, the splittings required for successful RL are generated radiatively [13] by renormalization group (RG) evolution and can therefore be naturally comparable to the decay widths of the heavy neutrinos.

τ-LEPTON FLAVOUR SYMMETRY

The leptonic Yukawa and Majorana sectors of the SM symmetrically extended with one singlet right-handed neutrino ν_{iR} per i family (with $i = 1,2,3 = e,\mu,\tau$) are given by the Lagrangian

$$-\mathscr{L}_{Y,M} \;=\; \bar{L}\Phi\mathbf{h}^{\ell}\, l_R \;+\; \bar{L}\tilde{\Phi}\mathbf{h}^{\nu}\, \nu_R \;+\; \overline{\nu}_R^{C}\,\mathbf{m}_M\, \nu_R \;+\; \text{H.c.}, \tag{1}$$

where Φ is the SM Higgs doublet and $\tilde{\Phi} = i\tau_2\Phi^*$ its isospin conjugate. Moreover, we have suppressed the generation index i from the left-handed doublets $L_i = (\nu_{iL}, l_{iL})^T$, the right-handed charged leptons l_{iR} and the right-handed neutrinos ν_{iR}, while ordinary multiplication between vectors and matrices is implied.

TABLE 1. *Flavour charge assignments for the breaking* $SO(3) \to U(1)_{L_e+L_\mu} \times U(1)_{L_\tau}$.

	L_e, e_R	L_μ, μ_R	L_τ, τ_R	ν_1	$\nu_2 \pm i\nu_3$
$U(1)_{L_e+L_\mu}$	+1	+1	0	0	± 1
$U(1)_{L_\tau}$	0	0	+1	0	0

To obtain a phenomenologically relevant model in this minimal setup, at least 3 singlet heavy Majorana neutrinos $\nu_{1,2,3R}$ are needed and these have to be nearly degenerate in mass [10]. To ensure the latter, we assume that to leading order, the singlet Majorana sector is SO(3) symmetric, i.e.

$$\mathbf{m}_M \;=\; m_N \mathbf{1}_3 \;+\; \Delta\mathbf{m}_M \;, \tag{2}$$

where $\mathbf{1}_3$ is the 3×3 identity matrix and $\Delta\mathbf{m}_M$ is a general SO(3)-breaking matrix induced by RG effects. As we will discuss below, compatibility with the observed light neutrino masses and mixings requires that $(\Delta\mathbf{m}_M)_{ij}/m_N \lesssim 10^{-7}$, for electroweak-mass Majorana neutrinos, i.e. for $m_N \approx 0.1$–1 TeV. We will explicitly demonstrate, how such an SO(3)-breaking matrix $\Delta\mathbf{m}_M$, of the required order, can be generated radiatively via the RG evolution of the right-handed neutrino mass matrix $\mathbf{m}_M$ from the GUT scale $M_X \approx 2\times 10^{16}$ GeV to the mass scale of the right-handed neutrinos m_N.

In the physical charged-lepton mass basis, the SO(3) symmetry imposed on the singlet Majorana sector at the GUT scale M_X gets explicitly broken by a set of neutrino Yukawa couplings to the subgroup of lepton symmetries: $U(1)_{L_e+L_\mu} \times U(1)_{L_\tau}$. The flavour charge assignments that give rise to such a breaking are presented in Table 1.

As a consequence of the $U(1)_{L_e+L_\mu} \times U(1)_{L_\tau}$ symmetry, the neutrino Yukawa coupling matrix takes on the general form:

$$\mathbf{h}^\nu \;=\; \begin{pmatrix} 0 & ae^{-i\pi/4} & ae^{i\pi/4} \\ 0 & be^{-i\pi/4} & be^{i\pi/4} \\ 0 & 0 & 0 \end{pmatrix} + \begin{pmatrix} \varepsilon_e & 0 & 0 \\ \varepsilon_\mu & 0 & 0 \\ \varepsilon_\tau & \kappa_1 e^{-i(\pi/4-\gamma_1)} & \kappa_2 e^{i(\pi/4-\gamma_2)} \end{pmatrix}, \tag{3}$$

where the second term vanishes, if the symmetry is exact. In this symmetric limit, the light neutrinos remain massless to all orders in perturbation theory, whilst a and b are free unconstrained complex parameters. The phases accompanying these parameters in (3) are simply chosen for convenience to maximize the lepton asymmetry in leptogenesis when a and b are real. In order to give masses to the light neutrinos, the second term describes the departure from the flavour symmetric limit with $|\varepsilon_{e,\mu,\tau}|, \kappa_{1,2} \ll |a|, |b|$ and the phases $\gamma_{1,2}$ are unrestricted.

It is important to notice that the flavour structure of the neutrino Yukawa couplings $\mathbf{h}^\nu$ is preserved through the RG evolution, as long as the flavour symmetry is only weakly broken. In detail, RG effects violate the SO(3)-invariant form of the Majorana mass matrix $\mathbf{m}_M(M_X) = m_N \mathbf{1}_3$ by the 3-by-3 matrix $\Delta\mathbf{m}_M$. In the leading-log approximation, the SO(3)-breaking matrix $\Delta\mathbf{m}_M$ receives the following radiative corrections:

$$\Delta\mathbf{m}_M = -\frac{m_N}{8\pi^2} \ln\left(\frac{M_X}{m_N}\right) \mathrm{Re}\left[\mathbf{h}^{\nu\dagger}(M_X)\mathbf{h}^\nu(M_X)\right]. \tag{4}$$

Given the form invariance of $\mathbf{h}^\nu$ under RG effects, we may therefore define all input parameters at the right-handed neutrino mass scale m_N, where the matching with the light neutrino data is performed.

We may now determine the Yukawa parameters $(a, b, \varepsilon_e, \varepsilon_\mu, \varepsilon_\tau)$, in terms of the light-neutrino mass matrix $\mathbf{m}^\nu$ in the positive and diagonal charged lepton Yukawa basis. To do so, we first notice that the chosen symmetry $U(1)_{L_e+L_\mu} \times U(1)_{L_\tau}$ ensures the vanishing of the light neutrino mass matrix $\mathbf{m}^\nu$. However, the symmetry-breaking parameters induce a non-zero light neutrino mass matrix $\mathbf{m}^\nu$, which to leading order in these parameters is given by [11, 12]

$$\begin{aligned} \mathbf{m}^\nu \;&=\; -\frac{v^2}{2}\mathbf{h}^\nu \mathbf{m}_M^{-1} \mathbf{h}^{\nu\mathrm{T}} \;=\; \frac{v^2}{2m_N}\left(\frac{\mathbf{h}^\nu \Delta\mathbf{m}_M \mathbf{h}^{\nu\mathrm{T}}}{m_N} \;-\; \mathbf{h}^\nu\mathbf{h}^{\nu\mathrm{T}}\right) \\ &=\; -\frac{v^2}{2m_N}\begin{pmatrix} \frac{\Delta m_N}{m_N}a^2 - \varepsilon_e^2 & \frac{\Delta m_N}{m_N}ab - \varepsilon_e\varepsilon_\mu & -\varepsilon_e\varepsilon_\tau \\ \frac{\Delta m_N}{m_N}ab - \varepsilon_e\varepsilon_\mu & \frac{\Delta m_N}{m_N}b^2 - \varepsilon_\mu^2 & -\varepsilon_\mu\varepsilon_\tau \\ -\varepsilon_e\varepsilon_\tau & -\varepsilon_\mu\varepsilon_\tau & -\varepsilon_\tau^2 \end{pmatrix}, \end{aligned} \tag{5}$$

where $v = 2M_W/g_w = 245$ GeV is the vacuum expectation value of the SM Higgs field Φ. In deriving the last equation in (5), we have also assumed that $\sqrt{\frac{\Delta m_N}{m_N}}\,\kappa_{1,2} \ll \varepsilon_{e,\mu,\tau}$, where Δm_N stands for the expression

$$\Delta m_N \equiv 2(\Delta\mathbf{m}_M)_{23} + i\Big[(\Delta\mathbf{m}_M)_{33} - (\Delta\mathbf{m}_M)_{22}\Big] = -\,\frac{m_N}{8\pi^2}\ln\left(\frac{M_X}{m_N}\right)\Big[2\kappa_1\kappa_2\sin(\gamma_1+\gamma_2) + i(\kappa_2^2-\kappa_1^2)\Big]\,. \qquad (6)$$

As a consequence of the flavour symmetry $\mathrm{U}(1)_{L_e+L_\mu}\times\mathrm{U}(1)_{L_\tau}$, the symmetry-violating parameters $\varepsilon_{e,\mu,\tau}$ and $\kappa_{1,2}$ enter the tree-level light neutrino mass matrix $\mathbf{m}^\nu$ in (5) quadratically. This in turn implies that for electroweak-scale heavy neutrinos $m_N \sim v$, the symmetry-breaking Yukawa couplings need not be much smaller than the electron Yukawa coupling $h_e \sim 10^{-6}$. Moreover, one should observe that only a particular combination of SO(3)-violating terms $(\Delta\mathbf{m}_M)_{ij}$ appears in $\mathbf{m}^\nu$ through Δm_N. Nevertheless, for electroweak-scale heavy neutrinos with mass differences $|\Delta m_N|/m_N \lesssim 10^{-7}$, one should have $|a|,\,|b| \lesssim 10^{-2}$ to avoid getting too large light neutrino masses much above 0.5 eV. Finally, we note that one-loop Z- and Higgs-boson effects can induce additional contributions to the light neutrino mass matrix $\mathbf{m}^\nu$ [14]. As well as being loop suppressed, these radiative contributions are proportional to the symmetry-breaking parameters and can therefore be neglected for the RτL scenario under study [11].

Given the analytic form (5) of the light neutrino mass matrix, we may directly compute the neutrino Yukawa couplings $(a, b, \varepsilon_e, \varepsilon_\mu, \varepsilon_\tau)$, as functions of m_N, the phenomenologically constrained neutrino mass matrix $\mathbf{m}^\nu$ and the symmetry-breaking parameters $\kappa_{1,2}$ and $\gamma_{1,2}$:

$$\begin{aligned}
a^2 &= \frac{2m_N}{v^2}\,\frac{8\pi^2}{\ln(M_X/m_N)}\left(m^\nu_{11} - \frac{(m^\nu_{13})^2}{m^\nu_{33}}\right)\Big[2\kappa_1\kappa_2\sin(\gamma_1+\gamma_2) + i(\kappa_2^2-\kappa_1^2)\Big]^{-1}\,,\\
b^2 &= \frac{2m_N}{v^2}\,\frac{8\pi^2}{\ln(M_X/m_N)}\left(m^\nu_{22} - \frac{(m^\nu_{23})^2}{m^\nu_{33}}\right)\Big[2\kappa_1\kappa_2\sin(\gamma_1+\gamma_2) + i(\kappa_2^2-\kappa_1^2)\Big]^{-1}\,,\\
\varepsilon_e^2 &= \frac{2m_N}{v^2}\,\frac{(m^\nu_{13})^2}{m^\nu_{33}}\,,\qquad \varepsilon_\mu^2 = \frac{2m_N}{v^2}\,\frac{(m^\nu_{23})^2}{m^\nu_{33}}\,,\qquad \varepsilon_\tau^2 = \frac{2m_N}{v^2}\,m^\nu_{33}\,.
\end{aligned} \qquad (7)$$

Since the approximate light-neutrino mass matrix $\mathbf{m}^\nu$ in (5) has rank 2, the lightest neutrino mass eigenstate ν_1 will be massless in this approximation. The relations given in (7) will be used to obtain numerical estimates of the BAU, in terms of m_N and the symmetry-breaking parameters $\kappa_{1,2}$ and $\gamma_{1,2}$, for both normal and inverted hierarchy scenarios of light neutrinos.

RESONANT LEPTOGENESIS

Within the framework of leptogenesis, a net non-zero leptonic asymmetry results from the CP-violating decays of the heavy Majorana neutrinos N_α into the left-handed charged leptons l^-_L and light neutrinos ν_{lL}. Consequently, we have to calculate the partial decay width of the heavy Majorana neutrino N_α into a particular lepton flavour l,

$$\Gamma_{\alpha l} \;=\; \Gamma(N_\alpha \to l^-_L + W^+) + \Gamma(N_\alpha \to \nu_{lL} + Z, H)\,. \qquad (8)$$

In RL models, resumming the absorptive parts of the heavy Majorana-neutrino self-energy transitions $N_\beta \to N_\alpha$ plays an important role in the computation of $\Gamma_{\alpha l}$ [5, 6].

The leptonic asymmetries for each individual lepton flavour can be expressed in terms of the resummed neutrino Yukawa couplings $\overline{\mathbf{h}}^\nu_{l\alpha}$ [6, 15] as

$$\delta_{\alpha l} \;\equiv\; \frac{\Gamma_{\alpha l} - \Gamma^C_{\alpha l}}{\sum\limits_{l=e,\mu,\tau}\Big(\Gamma_{\alpha l} + \Gamma^C_{\alpha l}\Big)} \;=\; \frac{\big|\overline{\mathbf{h}}^\nu_{l\alpha}\big|^2 - \big|\overline{\mathbf{h}}^{\nu C}_{l\alpha}\big|^2}{\big(\overline{\mathbf{h}}^{\nu\dagger}\overline{\mathbf{h}}^\nu\big)_{\alpha\alpha} + \big(\overline{\mathbf{h}}^{\nu C\dagger}\overline{\mathbf{h}}^{\nu C}\big)_{\alpha\alpha}}\,. \qquad (9)$$

The analytic results for the leptonic asymmetries $\delta_{\alpha l}$ simplify considerably in the 2-heavy neutrino mixing limit [5, 6],

$$\delta_{\alpha l} \;\approx\; \frac{\mathrm{Im}\big[(\mathbf{h}^{\nu\dagger}_{\alpha l}\mathbf{h}^\nu_{l\beta})(\mathbf{h}^{\nu\dagger}\mathbf{h}^\nu)_{\alpha\beta}\big]}{(\mathbf{h}^{\nu\dagger}\mathbf{h}^\nu)_{\alpha\alpha}(\mathbf{h}^{\nu\dagger}\mathbf{h}^\nu)_{\beta\beta}}\;\frac{(m^2_{N_\alpha} - m^2_{N_\beta})\,m_{N_\alpha}\Gamma^{(0)}_{N_\beta}}{(m^2_{N_\alpha} - m^2_{N_\beta})^2 + m^2_{N_\alpha}\Gamma^{(0)2}_{N_\beta}}\,, \qquad (10)$$

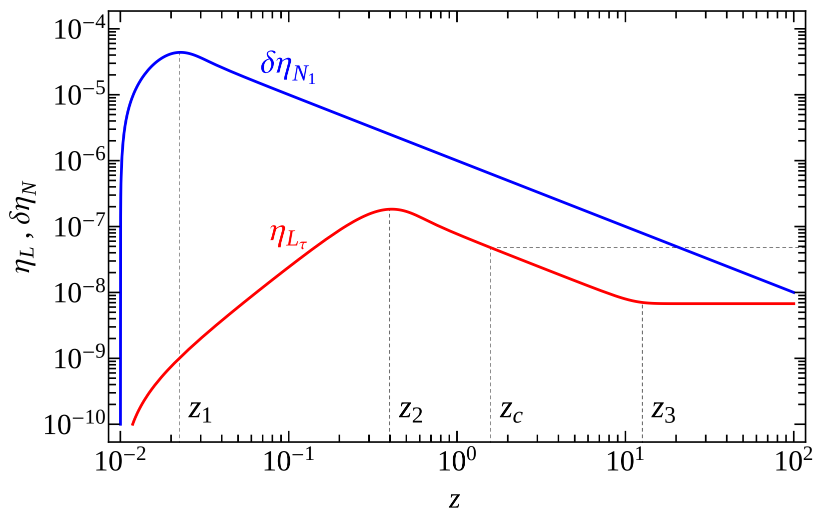

FIGURE 1. Numerical solutions to the Boltzmann equations for $\delta\eta_{N_1} = \eta_{N_1}/\eta_N^{\rm eq} - 1$ and η_{L_τ}, respectively, in a RτL model with $m_N = 220$ GeV, $K_{N_1} = \Gamma_{N_1}/H_N = 10^6$, $K_\tau = 10^2$, $\delta_\tau = 10^{-6}$.

where $\alpha, \beta = 1,2$ and $\Gamma^{(0)}_{N_a}$ is the tree-level decay width of the heavy Majorana neutrino N_α. The following two conditions for having resonantly enhanced leptonic asymmetries $\delta_{\alpha l} \sim \mathscr{O}(1)$ may be derived [5, 11]:

$$\text{(i)} \qquad |m_{N_\alpha} - m_{N_\beta}| \;\sim\; \frac{\Gamma_{N_{\alpha,\beta}}}{2}\,, \tag{11}$$

$$\text{(ii)} \qquad \frac{\left|\mathrm{Im}\left[\,(\mathbf{h}^{\nu\dagger}_{\alpha l}\mathbf{h}^{\nu}_{l\beta})\,(\mathbf{h}^{\nu\dagger}\mathbf{h}^{\nu})_{\alpha\beta}\right]\right|}{(\mathbf{h}^{\nu\dagger}\mathbf{h}^{\nu})_{\alpha\alpha}\,(\mathbf{h}^{\nu\dagger}\mathbf{h}^{\nu})_{\beta\beta}} \;\sim\; 1\,. \tag{12}$$

Within our RL scenarios, the first condition in (11) is naturally fulfilled as the heavy-neutrino mass splittings are generated via RG effects and are of the required order. The second condition is crucial as well and controls the size of the leptonic asymmetries. A detailed comparison of the predictions obtained by different approaches to resonant leptogenesis is given in [12].

The Boltzmann equations for the heavy Majorana neutrino and lepton-number densities, $\eta_{N_{1,2,3}}$, $\eta_{L_{e,\mu,\tau}}$, normalized to the photon number density n_γ are given by [11, 12]

$$\frac{d\eta_{N_\alpha}}{dz} \;=\; -\,\frac{z}{n_\gamma H_N}\left(\frac{\eta_{N_\alpha}}{\eta_N^{\rm eq}} - 1\right)\gamma^{N_\alpha}_{L\Phi}\,, \tag{13}$$

$$\begin{aligned}\frac{d\eta_{L_l}}{dz} \;=\; & \frac{z}{n_\gamma H_N}\left[\,\sum_{\alpha=1}^{3}\left(\frac{\eta_{N_\alpha}}{\eta_N^{\rm eq}} - 1\right)\delta_{\alpha l}\gamma^{N_\alpha}_{L\Phi} \;-\; \frac{2}{3}\,\eta_{L_l}\sum_{k=e,\mu,\tau}\left(\gamma^{L_l\Phi}_{L_k^C\Phi^\dagger} + \gamma^{L_l\Phi}_{L_k\Phi}\right)\right.\\ & \left. -\,\frac{2}{3}\sum_{k=e,\mu,\tau}\eta_{L_k}\left(\gamma^{L_k\Phi}_{L_l^C\Phi^\dagger} - \gamma^{L_k\Phi}_{L_l\Phi}\right)\right]\,, \end{aligned} \tag{14}$$

where $\alpha = 1,2,3$ and $l = e,\mu,\tau$. In addition, $H_N \approx 17 \times m_N^2/M_{\rm P}$ is the Hubble parameter at $T = m_N$, where $M_{\rm P} = 1.2 \times 10^{19}$ GeV is the Planck mass. The T-dependence of the BEs (13) and (14) is expressed in terms of the dimensionless parameter $z = \frac{m_N}{T}$, and $\eta_N^{\rm eq}$ is the equilibrium number density of the heavy Majorana neutrino N_α, normalized to the number density of photons. The BEs (13) and (14) include the collision terms $\gamma(X \to Y) \equiv \gamma^X_Y$ for the decays $N_\alpha \to L_l\Phi$, as well as the $\Delta L = 0,2$ resonant scattering processes $L_k\Phi \to L_l\Phi$ and $L_k\Phi \to L_l^C\Phi^\dagger$ [11].

The lepton asymmetries are then partially converted via $(B+L)$-violating sphaleron processes [16]. For temperatures T larger than the critical temperature $T_c \approx 135$ GeV of the electroweak phase transition, the conversion of the total lepton-to-photon ratio $\sum_{l=e,\mu,\tau}\eta_{L_l}$ to the baryon-to-photon ratio η_B^c at T_c is given by $\eta_B^c = -\frac{28}{51}\sum_{l=e,\mu,\tau}\eta_{L_l}$.

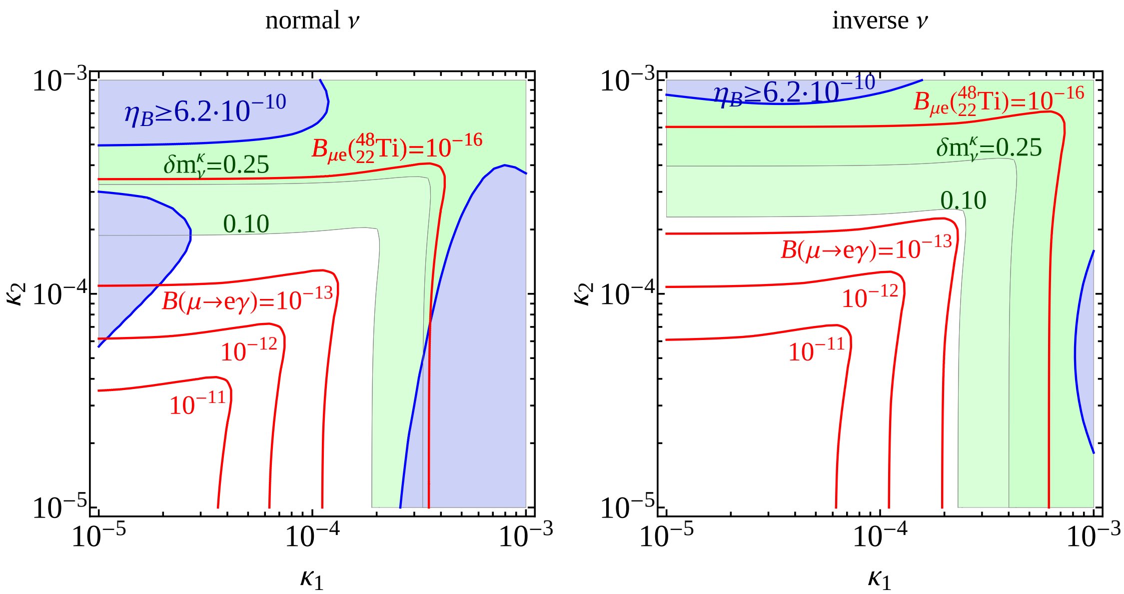

FIGURE 2. Baryon asymmetry η_B (blue contours) and the LFV observables $B(\mu \to e\gamma)$ and $B_{\mu e}(^{48}_{22}\mathrm{Ti})$ (red contours) as functions of κ_1 and κ_2 in the RτL model with $m_N = 120$ GeV, assuming a normal (left) and inverse (right) light neutrino mass spectrum. The remaining parameters are: $\gamma_1 = 3\pi/8$, $\gamma_2 = \pi/2$, $\phi_1 = \pi$, $\phi_2 = 0$, $\mathrm{Re}(a) > 0$. The neutrino oscillation parameters are set at their best fit values, with $\sin^2\theta_{13} = 0.027$. The blue shaded regions denote the parameter space where the baryon asymmetry is larger than the observed value. The green shaded areas labeled as '$\delta m_\nu^\kappa = 0.25$' and '0.10' indicate the parameter space where the inversion of the light-neutrino mass matrix is violated at the 25% and 10% level, respectively.

After dilution of this baryon asymmetry η_B^c by standard photon interactions until the recombination epoch, $\eta_B \approx 1/27\eta_B^c$, the theoretical prediction can now be compared with the current observational value for the baryon-to-photon asymmetry [2]: $\eta_B^{\mathrm{obs}} = (6.20 \pm 0.15) \times 10^{-10}$.

In the limit of strong washout, the baryon asymmetry η_B becomes relatively independent of the initial values of η_{L_l} and η_{N_α}. In this case, the Boltzmann equations can be solved analytically and the resulting BAU can be estimated to be [12]

$$\eta_B \approx \frac{-3 \cdot 10^{-2}\delta_\tau}{K_\tau \min\left[m_N/T_c, 1.25\ln(25K_\tau)\right]}, \tag{15}$$

with the total τ asymmetry $\delta_\tau = \sum_\alpha \delta_{\alpha\tau}$ and the total τ asymmetry washout factor $K_\tau = \sum_\alpha \Gamma_{N_\alpha \to L_\tau \phi}/H_N$. This formula provides a good approximation of the generated BAU in the RτL scenario within the regime $K_\tau \gtrsim 5$ for a right-handed neutrino mass scale of the order of the EW scale. Hence, to account for the observed BAU, a τ lepton asymmetry $\delta_\tau \gtrsim 10^{-7}$ is required.

LEPTON FLAVOUR VIOLATION AND θ_{13}

Heavy Majorana-neutrino loop effects may induce sizeable LFV couplings to the photon and the Z boson. These couplings give rise to LFV decays, such as $\mu \to e\gamma$ [17], $\mu \to eee$ [18] and $\mu \to e$ conversion in nuclei. In the RτL model, only two of the right-handed neutrinos have appreciable e- and μ-Yukawa couplings $a, b \approx 10^{-2}$ for $\kappa_{1,2} \approx 10^{-5} - 10^{-3}$. The LFV branching ratio $B(\mu \to e\gamma)$ can be expressed as [12]

$$B(\mu \to e\gamma) \approx 8.0 \cdot 10^{-4} \times g\left(\frac{m_N}{m_W}\right)\frac{v^4}{m_N^4}a^2b^2, \tag{16}$$

with the loop function $g(x)$, possessing the limits $g \to 1$ for $m_N \gg m_W$ and $g \to 1/16$ for $m_N = m_W$. This theoretical prediction can be favourably contrasted with the current experimental upper limit $B(\mu \to e\gamma) < 2.4 \cdot 10^{-12}$ [19].

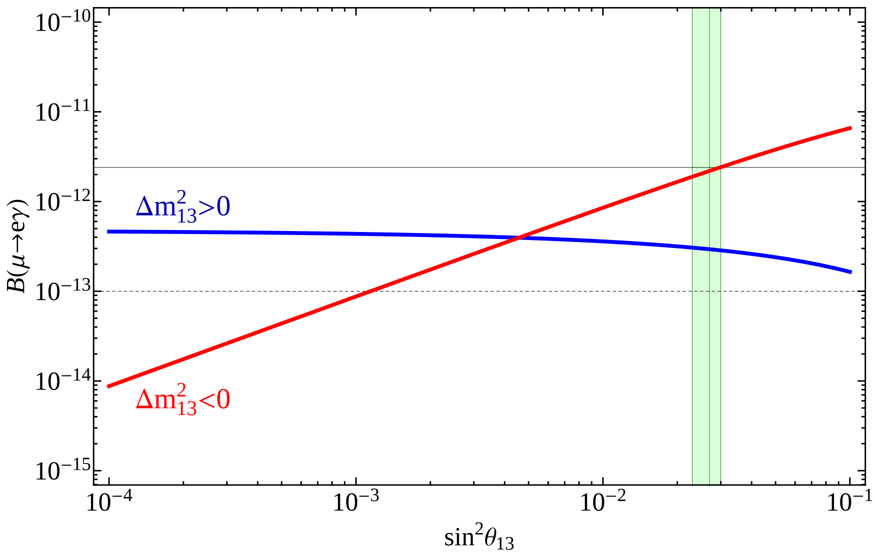

FIGURE 3. Branching ratio $B(\mu \to e\gamma)$ as a function of $\sin^2\theta_{13}$, in RτL scenarios with a normal (blue) and inverted (red) light neutrino mass spectra. The model parameters for these two cases are as in Fig. 2, with $\kappa_1 = 10^{-5}$ and $\kappa_2 = 10^{-4}$. The horizontal solid (dashed) line denotes the current (expected future) limit on $B(\mu \to e\gamma)$. The vertical green band indicates the allowed range of $\sin^2\theta_{13}$, with an uncertainty of $\pm 1\sigma$; the middle vertical line shows the current experimental best fit value for $\sin^2\theta_{13}$ [20].

As discussed above, the neutrino oscillation data can be used to determine the theoretical parameters of the RτL model by inverting the seesaw formula and solving for the neutrino Yukawa couplings a, b and $\varepsilon_{e,\mu,\tau}$. In addition, the electroweak-scale flavour structure of the right-handed neutrino mass matrix $\mathbf{m}_M$ is generated from a flavour-universal heavy-neutrino mass matrix $\mathbf{m}_M(M_X) = m_N \mathbf{1}_3$ at the GUT scale M_X, after taking into account RG-running effects. We use the best fit values of [20] for the measured neutrino oscillation parameters, and the three unknown CP phases of the light neutrino sector are taken into account by scanning over possible combinations of the extreme CP parities.

Fig. 2 shows the numerical estimates of the baryon asymmetry η_B and the LFV observables $B(\mu \to e\gamma)$ and $B_{\mu e}({}^{48}_{22}\mathrm{Ti})$, as functions of the Yukawa-coupling parameters κ_1 and κ_2, in a RτL model with $m_N = 120$ GeV and with either a normal or an inverted light neutrino mass spectrum. The phases of the $\kappa_{1,2}$ parameters are: $\gamma_1 = 3\pi/8$, $\gamma_2 = 1/2\pi$. This choice approximately enhances the overlap of the areas with successful baryogenesis and large LFV rates. The blue shaded areas denote the parameter space where the numerically predicted baryon asymmetry η_B is larger than the observational value $\eta_B^{\mathrm{obs}} = 6.2 \times 10^{-10}$. These areas should be regarded as representing regions of viable parameter space for successful leptogenesis, given the freedom of re-adjusting the CP phases $\gamma_{1,2}$ of the Yukawa couplings $\kappa_{1,2}$. The green areas denote the parameter space where the assumed approximation $\kappa_{1,2} \ll a, b$ is violated at 10% or 25%.

The branching ratio $B(\mu \to e\gamma)$ is enhanced in the case of an inverted hierarchical light neutrino mass spectrum ($\Delta m^2_{13} < 0$). Instead, in the normal hierarchical light-neutrino scenario with $\Delta m^2_{13} > 0$, $B(\mu \to e\gamma)$ is essentially independent of $\sin\theta_{13}$. This is demonstrated in Fig. 3, showing $B(\mu \to e\gamma)$ as a function of $\sin^2\theta_{13}$, where $\kappa_1 = 10^{-5}$ and $\kappa_2 = 10^{-4}$. The same figure shows the allowed area (displayed by a vertical green band) for $\sin^2\theta_{13}$ [20], as deduced from the two recent experiments conducted by the Daya Bay and RENO collaborations [21, 22]. Unlike in the normal light-neutrino scenario, the predicted value for η_B in the RτL model with an inverted hierarchical light neutrino mass spectrum falls short of explaining the BAU by two orders of magnitude, for a potentially observable branching ratio of $B(\mu \to e\gamma) \sim 10^{-13}$.

CONCLUSIONS

Minimal low-scale seesaw scenarios of resonant leptogenesis are a class of models with interesting implications for observables of charged LFV, such as $\mu \to e\gamma$ and $\mu \to e$ conversion in nuclei. Here we have discussed aspects of

the RτL realization of resonant leptogenesis, where the observed BAU originates from an individual τ lepton-number asymmetry which gets resonantly enhanced via the out-of-equilibrium decays of nearly degenerate heavy Majorana neutrinos. The required mass splittings are generated radiatively via RG effects with universal right-handed neutrino masses at the GUT scale.

Analogous to the RτL case discussed here, ReL and RμL scenarios with approximately protected e, μ asymmetries, respectively, may be constructed [12]. Heavy Majorana neutrinos in such RℓL scenarios can be as light as 100 GeV, whilst their couplings to two of the charged leptons may be large, so as to lead to LFV effects that could be tested by the MEG and the COMET/PRISM experiments. In the RτL model with a normal light neutrino mass hierarchy, there is a sizeable model parameter space with successful leptogenesis and large LFV process rates, with $B(\mu \to e\gamma) \approx 10^{-12}$. This prediction is largely independent of $\sin^2\theta_{13}$ and the other light neutrino oscillation parameters. On the other hand, in the RτL model with inversely hierarchical light neutrinos, $B(\mu \to e\gamma)$ is linearly proportional to $\sin^2\theta_{13}$, and can be enhanced by more than one order of magnitude compared to the normal hierarchy case for $\sin^2\theta_{13}$ close to its upper experimental 2σ limit. Unfortunately, the generated baryon asymmetry η_B is suppressed in this scenario, and to test the viable parameter space for successful leptogenesis would require an experiment for $\mu \to e$ conversion in nuclei which is sensitive to $B_{\mu e} \approx 10^{-17}$–$10^{-16}$.

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